

Inflationary predictions at small γ

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ABSTRACT: This paper explores single field inflation models with a constant, but arbitrary speed of sound c_s , obtained by deforming the kinetic energy terms to a Dirac-Born-Infeld form. Allowing $c_s < 1$ provides a simple parameterization of non-gaussianity. The dependence of inflationary observables on the parameter c_s is considered in the leading order slow roll approximation. The results show that in most cases the dependence is actually rather weak for the range of c_s allowed by existing bounds on non-gaussianity.

KEYWORDS: Inflation.

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1. Introduction

Models of k-inflation [1, 2] have been the subject of much recent interest. These models modify standard single-field inflation by allowing a more general form of the kinetic energy terms. An important feature of k-inflation is the alteration of the speed of propagation of disturbances in the inflaton field – the speed of sound

c_s . Models of inflation with canonical kinetic energy have $c_s = 1$, but with a more general form of kinetic energy this is no longer the case – the speed of sound is a field dependent quantity.

One particularly interesting example of this is a model which replaces the canonical kinetic energy by the Dirac-Born-Infeld form [3, 4], which has come to the fore in recent years in the context of D-brane inflation [5]–[7]. It is however also quite natural to view models of this kind in the general context of k-inflation [8, 9]. The DBI form of the kinetic energy terms involves a square root factor, $\gamma > 1$, reminiscent of the Lorentz factor of special relativity. Indeed, the square root is responsible for introducing a “speed limit” on the inflaton scalar. This facilitates a form of non-slow-roll inflation [4], on which most effort has focused. For such trajectories of the inflaton the Lorentz factor γ is large. However, it is also interesting to consider what quantitative effect allowing $\gamma > 1$ has in the slow roll regime. In the case of DBI models the speed of sound is expressed in terms of γ as $c_s = 1/\gamma$.

A particularly simple situation arises when the γ factor is constant [10], which means that the speed of sound is also constant, as in the canonical case, but no longer equal to the speed of light. This case can be considered as a leading approximation in an expansion of the field dependent speed of sound if it is assumed to vary slowly in the relevant region of field space. Some exactly solvable examples of this sort have recently been discussed in [10]–[13].

From a phenomenological perspective the interest in k-inflationary models stems from the fact that they provide a fairly simple way of accounting for non-gaussianity [14] in the spectrum of density perturbations. So far no conclusive evidence for departures from gaussianity has been observed, but it is of great interest to see if, and what kind of, theoretical possibilities exist, given the high expectations that data which will become available in the coming years will allow for a significant tightening of existing bounds. The simplest measure of non-gaussianity is the parameter f_{NL} , which is defined in terms of 3-point functions in [15, 16] (for example). For DBI models one has the simple relation [15] $f_{NL} = 35(\gamma^2 - 1)/108$, so the value of γ strongly affects the deviations from gaussianity. The current bounds on f_{NL} imply roughly $\gamma < 35$. DBI models with constant speed of sound provide a very simple

parameterization of non-gaussianity in models of single field inflation. Even if the observed non-gaussianity turns out to be large, the variation of the speed of sound with scale could be negligible.

This note reports the results of analyzing a series of frequently considered models of slow roll inflation and explores how sensitive their predictions are when one allows a small deviation from the canonical form of the kinetic energy, as measured by a constant $\gamma > 1$. The observables n_s and r have been computed for general models of k-inflation by Garriga and Mukhanov [2], and they are easily adapted to the case of DBI kinetic energy [15]. These formulae are evaluated at the time when the present Hubble scale crossed the horizon during inflation. The corresponding number of e-folds is denoted by N_\star ¹. This number is afflicted by some theoretical uncertainties [17], and the possibility of having $\gamma > 1$ adds another ones, as discussed in section (2). The remaining sections discuss a number of popular inflationary models case by case. In most cases, notably chaotic inflation, the results for the inflationary observables do not depend on γ , or the dependence is very weak. However in some cases of modular inflation it is found that the tensor fraction r effectively grows with γ , so one can envisage that it might become observable due to this effect.

2. Basic formalism

The simplest models assume that inflation is driven by a single scalar field, whose contribution to the energy density dominates and leads to the negative pressure which drives the accelerated expansion. Under very general assumptions [1, 15] the inflaton action takes the form

$$S = \frac{1}{2} \int \sqrt{-g} (R + P(X, \phi)) , \quad (2.1)$$

where $X = \frac{1}{2}(\partial\phi)^2$.

Brane inflation models in string theory [5]-[7] have attracted much attention to a specific example, the Dirac-Born-Infeld action:

$$S = - \int d^4x \, a(t)^3 \{ f(\phi)^{-1} (\sqrt{1 - f(\phi) \dot{\phi}^2} - 1) + V(\phi) \} . \quad (2.2)$$

¹As in a number of other studies, the number of e-folds is defined as decreasing to 0 at the end of inflation.

The function f appearing here can be related to the compactification geometry of the D-brane model. The action (2.2) leads to field equations for a perfect fluid

$$\dot{\rho} = -3H(p + \rho) , \quad (2.3)$$

$$3M_P^2 H^2 = \rho , \quad (2.4)$$

with

$$p = \frac{\gamma - 1}{f\gamma} - V(\phi) , \quad (2.5)$$

$$\rho = \frac{\gamma - 1}{f} + V(\phi) , \quad (2.6)$$

where

$$\gamma = \frac{1}{\sqrt{1 - f(\phi)\dot{\phi}^2}} . \quad (2.7)$$

It is convenient to express these equations in the Hamilton-Jacobi form [18, 19]

$$3M_P^2 H^2 - V = \frac{\gamma - 1}{f} . \quad (2.8)$$

where now

$$\gamma(\phi) = \sqrt{1 + 4M_P^4 f H'^2} . \quad (2.9)$$

To quantify the conditions under which inflation takes place one defines

$$\varepsilon = \frac{2M_P^2}{\gamma} \left(\frac{H'}{H} \right)^2 . \quad (2.10)$$

The condition for the Universe to be inflating is $\varepsilon < 1$. Just as in the canonical case, the leading order slow roll approximation entails dropping the H' dependence in (2.8), which means taking

$$3M_P^2 H^2 = V . \quad (2.11)$$

It is straightforward to write down corrections to observables related to the primordial perturbation spectra given the results of Garriga and Mukhanov [2]: the spectral indices can be written as follows:

$$n_s - 1 = -2\varepsilon + \eta + \sigma , \quad (2.12)$$

$$n_T = -2\varepsilon . \quad (2.13)$$

where

$$\eta = \frac{4M_P^2}{\gamma} \frac{H''}{H} - 2\varepsilon + \sigma , \quad (2.14)$$

$$\sigma = -\frac{2M_P^2}{\gamma} \frac{H'}{H} \frac{\gamma'}{\gamma} . \quad (2.15)$$

In the above formulae the right hand side is to be evaluated at horizon crossing, i.e. at $k = aH\gamma$. In the slow roll approximation one has

$$\varepsilon = \varepsilon_V , \quad (2.16)$$

$$\eta = 2\eta_V - 4\varepsilon_V , \quad (2.17)$$

where

$$\varepsilon_V = \frac{M_P^2}{2\gamma} \left(\frac{V'}{V} \right)^2 , \quad (2.18)$$

$$\eta_V = \frac{M_P^2}{\gamma} \frac{V''}{V} , \quad (2.19)$$

are the potential slow roll parameters. These have the usual dependence on the potential, but include a factor of $1/\gamma = 1$. For the case of constant γ one also has $\sigma = 0$ in (2.12), so that

$$n_s - 1 = -6\varepsilon_V + 2\eta_V , \quad (2.20)$$

$$n_T = -2\varepsilon_V . \quad (2.21)$$

The tensor to scalar ratio is [15]

$$r = \frac{16\varepsilon_V}{\gamma} . \quad (2.22)$$

One way to proceed is to express everything as a function of N – the number of e-folds. Using the convention that N decreases during inflation, reaching $N = 0$ at the end of inflation, one has $dN = -Hdt$, which leads to the formula

$$N(\phi) = -\frac{1}{\sqrt{2}} \int_{\phi}^{\phi_e} \frac{d\phi}{M_P} \sqrt{\frac{\gamma}{\varepsilon}} . \quad (2.23)$$

The value of N at horizon-crossing (denoted by N_*) can be determined as explained, for example, in [20, 21]. The essence is that N_* is given by

$$N_* = 50 + \delta N + \log \left(\frac{H_{end}}{H_*} \right) + \log(\gamma) . \quad (2.24)$$

where $\delta N = \pm 20$ accounts for various uncertainties in the post-inflationary evolution of the Universe [22] (such as the reheating temperature). The last two terms reflect features of inflation itself. The last term is a consequence of the change of speed of sound – horizon crossing is at $k = aH\gamma$. The next to last term accounts for the decrease in energy density during inflation. It was discussed in [20, 21]. It is easy to show that in the DBI case one has

$$\log \left(\frac{H_{end}}{H_\star} \right) = \int_0^{N_\star} dN \, \varepsilon(N) . \quad (2.25)$$

This has the same form as in the case of canonical kinetic energy [20] (but of course the ε appearing here contains the γ factor as in (2.10)). While this term is zero for de Sitter expansion, it could in principle give a non-negligible contribution to (2.26). In such a case (2.26) is an equation which has to be solved for N_\star . In the examples considered in this paper the resulting shift of N is usually not large in the allowed range of γ : it can be of order of a few at most, so it is smaller than the uncertainties in δN . It does however affect the window of reasonable values of N_\star which is usually considered [22, 17].

In the sequel we will shift the number at which we evaluate N_\star closer to the upper limit of estimated values of e-folds' number, i.e. we will use the formula:

$$N_\star = 60 + \log \left(\frac{H_{end}}{H_\star} \right) + \log(\gamma) , \quad (2.26)$$

where δN has been partially absorbed in a number 60.

3. Large field models

3.1 Chaotic inflation

Chaotic inflation [23] has become one of the most important examples of inflation. This is due partly to its simplicity – essentially any potential will work, and for simple monomial examples there is only one parameter, which can be fixed using the COBE normalization condition. For a quadratic monomial potential the model then predicts the scalar index, and the prediction is consistent with most recent data [28]. The simplicity of chaotic inflation is not to be held against it: one can envisage it

arising as an effective description in more complex contexts, such as supergravity or string theory [24]. The fact that chaotic inflation with a monomial potential gives a firm prediction for the scalar index makes it particularly interesting from the point of view of this note: it is natural to ask how sensitive this prediction is to deformations of the kinetic term.

The relevant potential reads:

$$V = \Lambda \left(\frac{\phi}{\mu} \right)^p, \quad (3.1)$$

In the slow roll approximation one has:

$$\varepsilon_V = \frac{M_P^2}{2\gamma} \left(\frac{V'}{V} \right)^2 \approx \frac{k^2}{2\gamma} \frac{M_P^2}{\phi^2}. \quad (3.2)$$

We need to express ε_H in terms of N and evaluate it at horizon-crossing. N is given in terms of ϕ by the formula (2.23)

$$N = \frac{\gamma}{2k} \left(\frac{\phi^2}{M_P^2} - \frac{\phi_e^2}{M_P^2} \right). \quad (3.3)$$

Inflation ends when $\varepsilon_H = 1$, i.e. when the inflaton field reaches the value $\frac{\phi_e^2}{M_P^2} = \frac{k^2}{2\gamma}$. Hence:

$$\frac{\phi^2}{M_P^2} = \frac{k}{\gamma} \left(2N + \frac{k}{2} \right). \quad (3.4)$$

This determines $\varepsilon_V(N)$ as

$$\varepsilon_V = \frac{k}{4N + k}. \quad (3.5)$$

The dependence on γ drops out in the above formula, so the scalar index does not explicitly depend on the value of γ . The fact is that $\gamma > 1$ only enters in by shifting the window of allowed values of N_* :

$$\log(H_i/H_e) = - \int_0^{N_*} dN \varepsilon_V(N). \quad (3.6)$$

This shift is quite small and has negligible effect on the observable quantities. Thus, at this level of approximation the predictions of chaotic inflation remain unchanged.

3.2 Natural inflation

Natural Inflation [26] is a model which arises from the assumption that the inflaton is an axion-like field, whose potential is generated by instanton effects:

$$V = \frac{1}{2} V_0 \left(1 - \cos \left(\sqrt{2|\eta_0|} \frac{\phi}{M_p} \right) \right). \quad (3.7)$$

Such a potential appears provides an example of a theoretically-motivated mechanism for generating chaotic inflation [27], since chaotic inflation in the case $k = 2$ may be considered as an approximation to a harmonic potential near its maximum.

The slow roll parameters read:

$$\varepsilon_V = \frac{|\eta_0|}{\gamma} \cot^2 \left(\sqrt{\frac{|\eta_0|}{2}} \frac{\phi}{M_p} \right) , \quad (3.8)$$

$$\eta_V = \varepsilon_V - \frac{|\eta_0|}{\gamma} . \quad (3.9)$$

To compute the inflationary observables one needs to assess the impact of $\gamma > 1$ on the allowed range of N_\star . Using (2.25) one finds

$$\ln H_i/H_e = -\frac{|\eta_0|}{\gamma} N_\star + \frac{1}{2} \ln \frac{\frac{|\eta_0| + \gamma}{\gamma} \exp \left(2 \frac{|\eta_0|}{\gamma} N_\star \right) - 1}{\frac{|\eta_0| + \gamma}{\gamma} - 1} , \quad (3.10)$$

which gives a shift of N_\star ranging from 62 (for $\gamma = 1$) to 66 (for $\gamma = 35$). This shift of N_\star lays well within the uncertainty due to unknown features of post-inflationary evolution (such as details of the reheating stage). Therefore, for the remaining calculations of natural inflation, the dependence of N_\star will be suppressed with the understanding that the allowed range is shifted by this amount.

Although N_\star depends weakly on γ , the spectral index $n_s - 1$, which is a physical observable (in contrast to N_\star) is a function of both γ and N_\star :

$$n_s - 1 = -6\varepsilon_V + 2\eta_V = -\frac{\frac{4|\eta_0|}{\gamma}}{\frac{|\eta_0| + \gamma}{\gamma} \exp \left(2 \frac{|\eta_0|}{\gamma} N_\star \right) - 1} - 2 \frac{|\eta_0|}{\gamma} . \quad (3.11)$$

This relation brings more significant change on the predicted by theory values of $n_s - 1$ – around 10% for $\gamma = 35$. The value of parameter $|\eta_0|$ was taken such as to mimic WMAP5 results [28] for $\gamma = 1$, i.e. $-\eta_0 = 0.014$. In this model, $n_s(\gamma)$ is an increasing function and tends to 1. However, the spectral tilt depends effectively on the ratio $|\eta_0|/\gamma$, so unless one has an independent estimate of $|\eta_0|$ in a specific model one can just fix this ratio using the observed value of $n_s - 1$.

For the tensor fraction r one has

$$r = \frac{16\varepsilon_V}{\gamma} = 16 \frac{\frac{|\eta_0|}{\gamma^2}}{\frac{|\eta_0| + \gamma}{\gamma} \exp\left(2\frac{|\eta_0|}{\gamma} N_\star\right) - 1} . \quad (3.12)$$

This depends on γ and $|\eta_0|/\gamma$ separately. If the value of $|\eta_0|/\gamma$ is fixed using the measured value of $n_s - 1$, then $r(\gamma) \sim 1/\gamma$. Taking into account the weak dependence of N_\star on γ (but with fixed $|\eta_0|/\gamma$) does not significantly alter the conclusions.

4. Small field models

4.1 New and modular inflation

In models of this type the inflaton field is usually assumed to be moving away from $\phi = 0$ and for some time the potential can be approximated by a polynomial of the form [22, 29, 30]:

$$V = V_0 \left(1 - \frac{\phi^k}{\mu^k}\right) . \quad (4.1)$$

This could be, for example, the Taylor expansion of potential arising from a phase transition associated with spontaneous symmetry breaking. This expansion is taken near unstable equilibrium at the origin, with k being the lowest non vanishing derivative there. It is assumed that $\phi \ll M_P$, and that the constant term dominates the potential.

To assess the impact of a having a lower speed of sound it is convenient to consider the cases $k = 2$ and $k \geq 3$ separately.

4.1.1 $k = 2$

In this case it is convenient to rewrite the inflation potential (4.1) in the form

$$V = V_0 \left(1 - \frac{1}{2} |\eta_0| \frac{\phi^2}{M_P^2}\right) , \quad (4.2)$$

where $\eta_0 < 0$ is the value of the parameter η_V at the maximum of the potential. Thus, in the slow roll approximation

$$\phi = \phi_e \exp\left(-\frac{|\eta_0|}{\gamma} N\right) , \quad (4.3)$$

and

$$\varepsilon_V = \frac{|\eta_0|^2}{2\gamma} \frac{\phi_e^2}{M_P^2} \exp\left(-\frac{2|\eta_0|}{\gamma} N\right) , \quad (4.4)$$

$$\eta_V = \frac{\eta_0}{\gamma} . \quad (4.5)$$

Here the value of the field ϕ_e at the end of inflation is kept as a parameter. This is due to the specific property of the scalar potential with a leading quadratic term [25, 27]. Namely, the small field assumption $\phi \ll M_P$, valid when the observable perturbations are generated, is not fulfilled during the whole inflation phase. The condition $\varepsilon = 1$ which is usually used to determine ϕ_e would in this case imply $\phi_e \gg M_P$. One expects however that the non-leading terms in the potential start playing an important role long before. Therefore, following [25, 27], we shall assume that due to these contributions the potential steepens and inflation ends when the inflation field ϕ is of order of the Planck's mass M_P .

Using the same procedure as in the previous section, the contribution to N_\star in (2.26) is

$$\ln H_i/H_e = \frac{|\eta_0|}{4} \frac{\phi_e^2}{M_P^2} \left[\exp\left(-\frac{2|\eta_0|}{\gamma} N_\star\right) - 1 \right] . \quad (4.6)$$

Since the spectral tilt $n_s - 1$ does not depend directly on N_\star :

$$n_s - 1 \approx 2\eta_H = -2\frac{|\eta_0|}{\gamma} . \quad (4.7)$$

(where we have taken into account $\varepsilon_V \ll \eta_V$), one can determine the ratio η_0/γ using the WMAP5 result for $n_s - 1$. One finds $\eta_0/\gamma = 0.02$. It can then be seen that the shift of N_\star is then no more than 4 efolds (allowing $1 \leq \gamma \leq 35$).

The tensor fraction r reads

$$r = 8\frac{|\eta_0|^2}{\gamma} \frac{\phi_e^2}{M_P^2} \exp\left(-\frac{2|\eta_0|}{\gamma} N_\star\right) = 2\frac{(n_s - 1)^2}{\gamma} \frac{\phi_e^2}{M_P^2} \exp((n_s - 1)N_\star) , \quad (4.8)$$

and for $\phi_e < M_P$ is unlikely to be observable. The small variation of N_\star with γ does not alter this situation.

4.1.2 $k = 3$ and higher

Here the results strongly depend on the value of the scale μ . For $\mu \ll M_p$ the spectral

tilt does not depend on γ : one finds

$$n_s - 1 = -\frac{k-1}{k-2} \frac{2}{N_\star}, \quad (4.9)$$

which is the standard result [27].

The shift of N_\star can be calculated as before and the result is

$$\ln H_i/H_e = \frac{1}{2} \left(\frac{k}{\sqrt{2}\gamma} \frac{M_p}{\mu} \right)^{\frac{k}{1-k}} \left\{ \left[(k-2) 2^{\frac{k-2}{2k-2}} k^{\frac{1}{k-1}} \gamma^{\frac{k}{2k-2}} \left(\frac{\mu}{M_p} \right)^{\frac{k}{1-k}} N_\star + 1 \right]^{\frac{k}{2-k}} - 1 \right\}. \quad (4.10)$$

This and the $\log \gamma$ term, give altogether N_\star changing with $\gamma \in [1, 35]$, for $k = 3$ from 60 to 65. This range gets narrower with decreasing ratio μ/M_p . However, as $n_s - 1 \sim 1/N_\star$, those modifications on N_\star do not have strong impact on values of spectral tilt.

If one allows μ being close to the Planck's mass both the spectral tilt and tensor fraction become quite sensitive to the value of γ . One finds

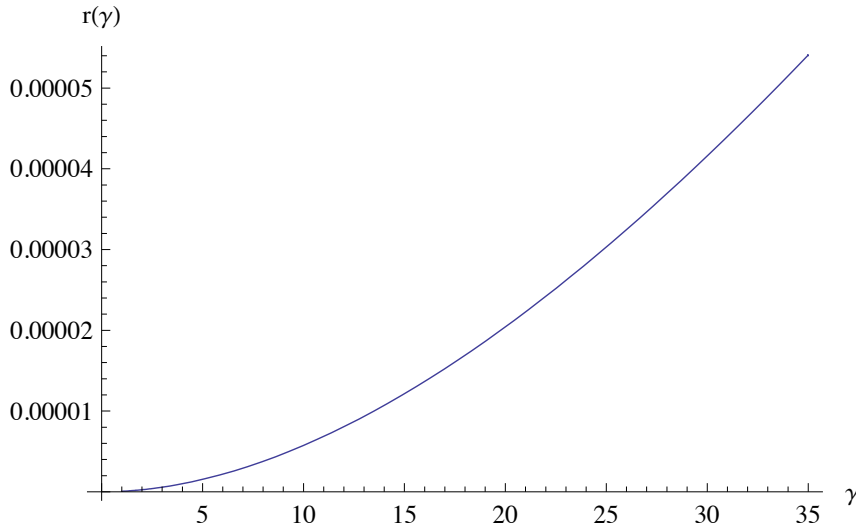
$$n_s - 1 = \frac{-2(k-1)}{(k-2)N_\star + k^{-\frac{1}{k-1}} \gamma^{\frac{k}{2k-2}} \left(\frac{\mu}{M_p} \right)^{\frac{k}{k-1}}}, \quad (4.11)$$

and

$$r = \frac{16}{\gamma} \left[(k-2) 2^{\frac{k-2}{2k-2}} k^{\frac{1}{k-1}} \gamma^{\frac{k}{2k-2}} \left(\frac{\mu}{M_p} \right)^{\frac{k}{1-k}} N_\star + 1 \right]^{\frac{2k-2}{2-k}}, \quad (4.12)$$

Increasing γ raises the value of n_s by a few percent and moves it closer to 1. For example, for $\mu \sim M_p$, $k = 3$ and $N_\star = 60$, $n_s - 1$ varies from -0.066 ($\gamma = 1$) to -0.058 ($\gamma = 35$). The tensor fraction r , for $k = 3$, $\gamma = 1$, $\mu \sim M_p$ and $N_\star = 60$ is of the order 10^{-7} , and even smaller for $\mu \approx 0.1 M_p$ – around 10^{-13} . However, the effect of increasing γ is quite significant, raising r to around 10^{-5} (when $\mu \sim M_p$), which is perhaps only an order of magnitude away from being observable [31]. This effect gets weaker for higher k , as well as for smaller μ/M_p .

The plot below shows $r(\gamma)$ for $k = 3$, $N_\star = 60$, $\mu \sim M_p$.



4.2 F - and D -term inflation

Another type of inflaton potential which is frequently considered appears in the context of supersymmetric models of inflation. For the simplest choice of Kähler potential, either F -term [32] or D -term [33] inflation is possible. The leading term in the potential is constant, and the dependence on the inflaton field arises as a loop correction. In many cases the relevant potential takes the form:

$$V(\phi) = V_0 \left(1 + \frac{\lambda^2}{8\pi^2} \ln \frac{\phi}{Q} \right) , \quad (4.13)$$

where λ is a coupling to the waterfall field and the renormalization scale Q is of the order of the inflationary field $\phi \gg \phi_e$.

In this case the predictions depend on γ very weakly. The number of e-folds N_\star does not depend on a Lorentz factor up to order λ^2 . The tensor fraction is a decreasing function of γ :

$$r = \frac{\lambda^2}{2\pi^2\gamma} \frac{1}{N_\star} \approx 0.0011 \frac{\lambda^2}{\gamma} \left(\frac{50}{N_\star} \right)^2 , \quad (4.14)$$

and for small λ much below 1, and bigger values of γ , values of r fall into unobservable region.

4.3 Exponential potential

The final model considered here is the exponential potential

$$V(\phi) = V_0 \left(1 - \exp \left(-q \frac{\phi}{M_p} \right) \right) , \quad (4.15)$$

which appears in the context of supergravity models (with the value $q = \sqrt{2}$) by choosing the Kähler potential appropriately ([33]). It may also arise (with $q = \sqrt{2/3}$) in models of inflation driven by non-Einstein gravity [34], although then, strictly speaking, one is on the border of the small-field regime.

In case of this potential, the γ -dependence of N_\star is negligible for either value of q , since

$$-\ln H_i/H_e = \frac{1/2}{\frac{q^2}{\gamma} N_\star + \frac{q + \sqrt{2}\gamma}{\sqrt{2}\gamma}} \approx \frac{\gamma}{2q^2 N_\star} < 1 . \quad (4.16)$$

However, the spectral index $n_s - 1$:

$$n_s - 1 \approx 2\eta_H = -\frac{2}{N_\star + \frac{\gamma}{q^2} + \sqrt{\frac{\gamma}{2q^2}}} , \quad (4.17)$$

shows rather insignificant dependence on γ : in supersymmetric model ($q = \sqrt{2}$), for a fixed $N_\star = 60$, $n_s - 1$ increases from -0.033 ($\gamma = 1$) to -0.032 for $\gamma = 35$. The situation is similar for the case of non-Einstein gravity inflation.

However, the tensor fraction r :

$$r(\gamma, N_\star) = \frac{8}{q^2 \left(N_\star + \frac{\gamma}{q^2} + \sqrt{\frac{\gamma}{2q^2}} \right)^2} , \quad (4.18)$$

is very sensitive to changing γ , and although $r(\gamma)$ is a decreasing function, it mainly remains in a range that is likely observable, i.e. it varies from 0.001 to 0.0006 with N_\star fixed at 60. Smaller values of N slightly increase the estimate of r .

5. Conclusions

It is expected that soon the existing bounds on non-gaussianity will be significantly tightened or a measurement of it will be made. In view of this it is important to

consider theoretical options which lead to non-gaussian perturbation spectra. One simple possibility was considered here. The conclusion is basically that as long as the speed of sound is constant, the basic predictions for the spectral tilt and tensor fraction are quite robust even if one allows f_{NL} as large as existing limits permit. Some variation of the tensor fraction with f_{NL} is possible however, and while generically r goes down with rising non-gaussianity, in some examples (new inflation with $k > 2$) it can increase somewhat.

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